

Universality of Glass Scaling in a $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ Thin Film

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The universal behavior of the continuous superconducting (glass) phase transition is studied in a $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ thin film. A new analysis technique for extracting critical exponents is developed, using current-voltage (J - E) measurements. The method places narrow limits on the glass scaling exponents, z_g and ν_g , as well as the glass transition temperature, $T_g(H)$. The universal values $z_g = 3.8 \pm 0.2$ and $\nu_g = 1.8 \pm 0.1$ are obtained. Using these values, the data are scaled for fields in the range 1-9 T. Collapse of the scaling forms at different fields describes the universal scaling function. The field dependence of the universal collapse is interpreted in terms of the 3D XY multicritical scaling theory.

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I. INTRODUCTION

The nature of the superconducting phase transition at finite fields in high temperature superconductors [HTS] remains surprisingly elusive. In clean systems, recent evidence suggests a first order melting transition of the vortex lattice.¹ However in disordered systems, the issue is not completely resolved: Does superconducting order exist at finite temperatures, separated from the normal state by a well-defined continuous phase transition? Or do manifestations of superconductivity at nonzero temperatures represent only nonequilibrium phenomena, as in the flux creep description?² This question is of great practical interest since the two mutually exclusive descriptions involve distinct dynamic equations which govern the response of driven vortices. In one case, the equations are mean-field-like;³ in the other, they describe a (glassy) critical state, which can occupy a very wide region of the superconducting phase diagram.⁴ From a fundamental viewpoint, the debate encompasses the nature of the superconducting state and long range order in the cuprate superconductors.

Shortly after the discovery of HTS, significant progress was made to clarify these points. Several types of glassy states⁵⁻⁷ were proposed as alternatives to the creep picture. A key feature of these glass phases is the true absence of resistance at low currents, for temperatures below the glass transition line, $T_g(H)$. Initial experiments strongly supported the presence of such a continuous phase transition,^{8,9} and the superiority of the glass description over the creep picture.¹⁰ However, further progress is hampered by the absence of a firm description of the glassy state, particularly regarding the nature of the glass transition. Neither the critical scaling exponents nor the scaling functions have been determined analytically, due to the difficulty of the theoretical problem. Further experimental treatments, while *individually* supporting a scaling description⁵ of current-voltage (J - E) characteristics, appear as a *group* to violate the notion of critical universality: both sample- and field-dependence

in the scaling exponents are observed by some authors.¹¹ A more rigorous test of universality, involving the scaling functions (which must also be universal according to the critical theory), has been investigated previously by Xenikos *et al.*¹² in terms of sample-dependence, by Wöltgens *et al.*¹³ in terms of the angle dependence of \mathbf{H} , and by Kötzler *et al.*¹⁴ in terms of magnetic field dependence. Since further theoretical advancements are not immediately forthcoming, confirmation of the existence of true finite temperature superconductivity now rests on the ability to discern universal behavior near the purported glass transition. We note, however, that this test provides evidence only for critical scaling, and gives limited insight into the nature of the glassy state.

Other types of critical fluctuations (besides glass fluctuations) may arise and compete with the glass excitations for critical dominance. If the scaling regions overlap, the corresponding scaling theories must be self-consistent, a fact which has a profound implications for universality. One such competing excitation in HTS is associated with the 3D XY model. The critical behavior of the zero field superconducting transition, T_c , is expected to belong to the 3D XY model class, in the strongly type-II limit.⁵ Furthermore, the finite field superconducting transition line $T_g(H)$ terminates at the 3D XY multicritical point: $T_c = T_g(H = 0)$. The topology of the H - T phase diagram therefore ensures that the glass and 3D XY critical regimes will overlap, and that their universal behaviors will be intimately related in some region about the 3D XY transition point.^{15,16} The cautions applied to any mean-field dynamic description in the glass scaling region also apply to the 3D XY region, which extends over a wide temperature range at low fields,^{17,18} encompassing the low field glassy region.^{15,19}

In this work, a single $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ thin film is studied. In order to explore universal behavior, a rigorous scaling treatment is required to place narrow limits on the glass scaling parameters. After restricting the analysis in this way, a J - E scaling procedure is performed, adjusting the scaling parameters within their bounds to obtain the best results. To check universality, glass scal-

ing is compared for different fields, not only in terms of scaling exponents, but also scaling functions; both are found to be field independent. However, fields do enter into the glass scaling in a way which is determined by the 3D XY multicritical theory at low fields. Although the present treatment is not applied to different samples here, it is hoped that such an extension may be performed in the near future. The following analysis is presented in the vortex glass formulation,⁵ however this does not confine the critical analysis in any way, since the different J - E scaling formulations are isomorphic to one another.

II. EXPERIMENT

A 4000 Å thick, deoxygenated $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ film, with the c -axis oriented perpendicular to the film surface, was used in this study. The film was prepared by pulsed laser deposition onto a heated (001) LaAlO_3 substrate, and initially had a superconducting transition temperature of $T_c = 92 \pm 2$ K at zero applied field. It was then deoxygenated at 400 °C using a controlled flow of oxygen and argon gasses, followed by a rapid quench to room temperature. The resulting superconductor had a transition temperature $T_c = 77$ K, corresponding to an oxygen stoichiometry of $\delta \simeq 0.24$. The film was patterned into a $120 \mu\text{m} \times 3100 \mu\text{m}$ bridge, using laser ablation.

The resistivity, $\rho(T)$, was measured as a function of temperature for fixed currents and fields. Measurements were then performed for a range of applied current values using a conventional four point geometry, for magnetic fields in the range $1 \leq H \leq 9$ T. The current densities ranged from 20 to 4×10^4 A cm⁻². The temperature range over which the resistivity was measured depended on the current. Typically, data were collected on warming for temperature windows of 25 K, starting at temperatures where the voltage signal increased over the instrumental noise level. This ensured that data were collected on both sides of the glass transition temperature, T_g . Throughout this work, errors are quoted as 2σ (standard deviation) unless otherwise noted.

III. SCALING ANALYSIS

A. Determination of z_g and T_g

The main focus of this work is the universal critical behavior near the glass transition temperature, T_g . Therefore a systematic, precise method for determining T_g , and the glass scaling exponents z_g and ν_g is required. One method for identifying these parameters involves scaling of current-voltage (J - E) isotherms, obtained in the vicinity of T_g for a single applied magnetic field (H).^{5,8} Typically,¹¹ the three scaling parameters are varied simultaneously to obtain the best collapse of the isotherms. With no constraints placed on the fitting parameters,

this method can obtain a reasonable scaling collapse for a large range of exponents. Conventionally, constraints are placed on the fitting procedure by identifying z_g and T_g from the critical isotherm,⁸ then requiring that the final scaling adjustments do not stray far from this estimate. However, low densities of isotherms, difficulties in subtracting the normal state background (causing the critical isotherm to appear as an imperfect power law), and Joule heating,⁸ all create ambiguities in the identification of the critical isotherm, reducing the precision of T_g to 1 K or worse.

To improve this situation, an alternative scaling procedure is developed below, in which z_g and ν_g are systematically obtained. This procedure places more rigorous constraints on the fitting parameters by largely circumventing the difficulties described above. An important feature of the method is its optimization of scaling in the vicinity of $T = T_g$; although this condition is required theoretically, it is not rigorously enforced in most scaling techniques. A similar technique was applied previously in the study of 3D XY , low field fluctuations in high- T_c superconducting films.¹⁵ In the present case, the starting point for the scaling procedure is the following ansatz:⁵

$$\sigma(J/T)^{(z_g-1)/2} = F \left[t(J/T)^{-1/2\nu_g} \right]. \quad (1)$$

Here, $\sigma = 1/\rho = J/E$ is the nonlinear conductivity, $F[x]$ is a scaling function, and $t = (T - T_g)/T_c$ is the relative temperature variable, for a fixed field H .

Since the left hand side of Eq. (1) displays no singularities, the function $F[x]$ must be smooth and single-valued. In particular, it is analytic at the glass transition $T = T_g$, where its argument $x = t(J/T)^{-1/2\nu_g}$ (the vortex glass scaling variable) is equal to 0. When $t = 0$, the scaling variable no longer depends on the current density J ; both sides of Eq. (1) should then be independent of J , provided that z_g is chosen correctly. The resistivity curves corresponding to different current values therefore all cross at T_g when plotted as in Fig. 1. In the inset, the exponents z_g , which provide the best simultaneous crossing, are shown for each applied field. The error bars reflect the fact that satisfactory crossings are obtained for some range of exponents. Since the field dependence of z_g is very weak, and since the error bars of z_g overlap over some range for all the fields studied here, the average value $z_g = 3.8 \pm 0.2$ is accepted as the result of this analysis. Note that for a given z_g , the corresponding value of T_g can be read off immediately from Fig. 1, with the small error of ± 0.02 K. However the total error in T_g , corresponding to the ambiguity in defining z_g , is found to be 0.4 K for the fields used here.

This *crossing point method* for identifying z_g and T_g optimizes scaling near $T = T_g$, and involves only the data taken near T_g . In the analysis developed above, the effects of the normal state conductivity have been ignored, assuming the fluctuation contributions to be much larger in comparison. To justify this assertion, the total non-linear conductivity is compared to the normal state

contribution at $T = T_g$, which is estimated roughly by extrapolating high temperature conductivity data. For the current densities used here, the total conductivity is at least 3 orders of magnitude larger than the normal state conductivity.

B. Determination of ν_g

Using the results of the previous section for T_g and z_g , the static scaling exponent ν_g can now be obtained through a full scaling of the conductivity data according to Eq. (1). However, it is beneficial to form an extension of the crossing point method of the previous section, which allows a value of ν_g to be identified for a specific (correlated) pair of parameters (z_g , T_g). This technique has the advantage of once again optimizing scaling near T_g . To augment the crossing point method, a second (and more conventional) method is also used to determine ν_g . This second technique complements the crossing point method by optimizing scaling in the ohmic (high-temperature) part of the scaling region. Of the two different methods, the first has greater theoretical justification, since it relies only on data in the vicinity of T_g . However, the two methods yield nearly identical results, with similar error levels. This supports the fact that the glass critical regime is large and robust.⁴ We now elaborate on the two methods.

Crossing point method, using data near $T = T_g$:

The method of the previous section can be extended to obtain terms involving ν_g . Since the conductivity scaling function $F[x]$ of Eq. (1) is analytic at $x = 0$, a Taylor series expansion may be performed about this point. For $t \simeq 0$, Eq. (1) becomes

$$\begin{aligned} \sigma(J/T)^{(z_g-1)/2} &= F[0] + t(T_g/J)^{1/2\nu_g} F'[0] + \dots \\ &= F[0] + mt + \dots \end{aligned} \quad (2)$$

Truncating Eq. (2) at its zeroth order term, $F[0]$, gives the z_g crossing point method used above. To determine ν_g , the first order term is also retained. The coefficient $m = (T_g/J)^{1/2\nu_g} F'[0]$ of the linear term in Eq. (2) is determined for different current densities J at $T = T_g$ by observing the slopes of the different curves in Fig. 1 at the crossing point. The result found above, $z_g = 3.8$, is used on the left hand side of Eq. (2). Since ν_g should be independent of the current, the values obtained in this way are averaged over the different currents to obtain error bars. The results of the crossing point method for ν_g are shown in Fig. 1 (inset) for each field. No obvious trend of ν_g with the field is apparent in this figure (and none is expected). It is therefore meaningful to average ν_g over field variations, giving a final result of $\nu_g = 1.77 \pm 0.03$. Keeping significant figures consistent with z_g , we quote $\nu_g = 1.8 \pm 0.1$.

Slope method, using data in the ohmic regime:

This makes use of the ohmic resistivity data in the limit of $J \rightarrow 0$ and $T > T_g$. Typically, data are obtained here

using current densities in the range 20 to 5×10^2 A cm⁻². In the ohmic regime, ρ is independent of J , so Eq. (1) reduces to $\rho \sim t^{\nu_g(z_g-1)}$. This power law behavior can be observed in Fig. 2. The slope of the linear (power law) portion of the curve can then be used to determine $\nu_g(z_g - 1)$, and subsequently ν_g , using $z_g = 3.8$. The apparent power law behavior of ρ , observed over two decades in Fig. 2, indicates that data used to obtain ν_g in this way remain within the critical scaling region. The resulting ν_g values are shown for different fields in the inset of Fig. 2. The error bars for each data point reflect the ambiguity in identifying the limits of the ohmic scaling region. Similar to the inset of Fig. 1, no obvious trend of ν_g with H is observed, nor is there any correlation between the field variations in the insets of Figs. 1 and 2. This provides additional support for identifying ν_g as a field average: $\nu_g = 1.79 \pm 0.04$, a result nearly identical to the crossing point method. Again, keeping significant figures consistent with z_g , we quote $\nu_g = 1.8 \pm 0.1$. This value is used in the following analysis.

IV. DISCUSSION OF UNIVERSALITY

The absence of field dependence in the exponents ν_g and z_g , demonstrated in Section III, is consistent with the vortex glass theory, which states that the critical behavior of the linear or nonlinear conductivity σ near the transition line $T_g(H)$ should belong to a single universality class, for some range of $H > 0$.⁵ The vortex glass excitations are thought to be formed within a medium of field-induced vortices, whose excursions from perfect rigid rod geometry, parallel to \mathbf{H} , are affected by entropy, applied currents, and the pinning effects of quenched disorder. The density of field vortices varies with H , while quenched disorder does not. For a particular sample then, vortex matter at different fixed fields forms unique vortex systems. The behavior and response of these unique systems to external probes may vary greatly with the field. Indeed, it has been noted that changes in the field may be responsible for phenomena such as dimensional crossover, causing vortex matter at low and high fields to behave very differently.²⁰ However, many of these phenomena are of mean-field origin, and should not be observed in measurements which predominantly reflect a critically diverging correlation length. It is not surprising then that vortex glass universality, observed here for all studied fields, $H > 0$, along $T_g(H)$ (and in Ref. 15 up to much higher fields) appears to hold in spite of such variations with field.

In order to observe universal glass behavior of both the critical exponents and the scaling function, it is necessary to introduce a new notation to account for field dependence in the vortex glass description. This field dependence is *not* specified by the vortex glass theory; it is not known to be universal, *a priori*, and may vary from sample to sample. The conductivity scaling of Eq. (1) can

then be rewritten as

$$\left(\frac{\rho}{\rho_0}\right) |t|^{-\nu_g(z_g-1)} = G_{\pm} \left[\frac{J}{T} \frac{T_g}{J_0} |t|^{-2\nu_g} \right]. \quad (3)$$

Here, we have introduced the field dependent resistivity scale $\rho_0(H)$ and the current scale $J_0(H)$. The transition line $T_g(H)$ is also field dependent. The new functions $G_{\pm}[x]$, corresponding to $T > T_g$ and $T < T_g$, and their arguments, are dimensionless.

To further develop a universal description, it is necessary to identify the scales $\rho_0(H)$ and $J_0(H)$ empirically. Methods for obtaining characteristic resistivity and current scales exist in the literature,⁸ however we believe the following definitions to be the most objective:

$$\rho_0 = \rho t^{-\nu_g(z_g-1)} \Big|_{T>T_g, J \rightarrow 0}, \quad (4)$$

$$J_0 = J(\rho/\rho_0)^{2/(1-z_g)} \Big|_{T=T_g}, \quad (5)$$

where ρ and J refer here to experimentally determined quantities. In this definition, ρ_0 is the ohmic resistivity scale, while J_0 is the current required to produce a resistance $\rho = \rho_0$ along the critical isotherm.

Using these definitions, the nonlinear resistivity curves, scaled at different fields according to Eq. (3), collapse as shown in Fig. 3. The choice of normalization in Eqs. (4) and (5) causes the dimensionless point (1, 1) to be located as shown in the figure. We emphasize that this collapse of J - E scaling curves, including different fields, demonstrates the universality of the vortex glass transition. By varying ν_g , z_g , and $T_g(H)$ within their error bars, as found above, it is not possible to obtain any discernible improvement in the scaling beyond Fig. 3. Additionally, for $z_g < 3.8$, the scaling deteriorates markedly.

The obtained field dependences of ρ_0 and J_0 are shown in Fig. 4. As mentioned above, neither of these quantities is specified by the vortex glass theory. However, they can be determined in a multicritical theory such as 3D XY , which is thought to apply to topological fluctuations at low fields.^{5,15,17} The following power law dependencies of $\rho_0(H)$ and $J_0(H)$ are derived from the 3D XY theory:¹⁹

$$\rho_0(H) = \bar{\rho} H^{[\nu_{xy}(z_{xy}-1) - \nu_g(z_g-1)]/2\nu_{xy}}, \quad (6)$$

$$J_0(H) = \bar{J} H^{1-(\nu_g/\nu_{xy})}, \quad (7)$$

where ν_{xy} and z_{xy} are 3D XY scaling exponents, and $\bar{\rho}$ and \bar{J} are unknown constants which are sample-, but not field-dependent.

In Fig. 4, $\rho_0(H)$ and $J_0(H)$ are described well by power laws for low fields.²¹ It is important to note that, in contrast to other 3D XY analyses performed in the literature, the observance of this power law does not hinge on an appropriate choice of T_c as a fitting parameter. In the fits shown, only data are used which fall within the expected 3D XY scaling region for this sample¹⁵ ($H < 8$ T). We find corresponding slopes of -3.1 ± 0.2 (ρ_0) and $-1.39 \pm 0.04 \simeq -1.4$ (J_0). A self-consistency

test is provided by Eq. (7), using the known value of $\nu_{xy} = 0.67$.²² The obtained value of $\nu_g = 1.60 \pm 0.03 \simeq 1.6$ provides a satisfactory check on the universality analysis. Additionally, z_{xy} may be calculated from Eq. (6), using the slope obtained in Fig. 4, together with $\nu_{xy} = 0.67$, $\nu_g = 1.8$, and $z_g = 3.8$. The result is $z_{xy} = 2.3 \pm 0.2$. [Note that the error sources in Eq. (6) are correlated.] This result is closer to that given by Booth *et al.*²³ than Moloni *et al.*¹⁵ We speculate that the difference between the present results and those of Ref. 15 may reflect the difficulty in treating background contributions to σ , which impedes the analysis of Ref. 15, but not the present one.

V. CONCLUSIONS

In this work, the vortex glass scaling parameters $\nu_g = 1.8 \pm 0.1$, $z_g = 3.8 \pm 0.2$, and $T_g(H)$ have been obtained through a systematic scaling technique, involving the vortex glass crossing point at $T = T_g$. Similar to other deoxygenated $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ films, the vortex glass scaling region was found to be wide, on the order of 10 K at $H = 7$ T. The absence of field dependence in the exponents ν_g and z_g provides a demonstration of universal behavior, since different fields correspond to unique vortex systems. A second (and stronger) demonstration of universality is given by the collapse of scaled J - E curves obtained at different fields in the vortex glass critical region, as shown in Fig. 3. Because universality is an essential component of critical phenomena, these checks on both the glass exponents and the scaling functions provide strong evidence for a critical phase transition line along $T_g(H)$ for $H > 0$. While additional evidence has not been presented here to demonstrate that this transition involves explicitly the broken symmetry associated with a vortex glass, it should be noted that the scaling functions in Fig. 3, below and above $T_g(H)$, exhibit the expected glassy characteristics at low currents: below $T_g(H)$ the non-linear resistivity decays exponentially to 0, while above $T_g(H)$ the resistivity decays exponentially to its normalized value of 1.

The values reported here for the dynamic scaling exponent, z_g , and the static scaling exponent, ν_g , fall within the range of previously reported values,¹¹ although our result $z_g = 3.8 \pm 0.2$ falls barely within the expected range⁵ $z_g \geq 4$. The significant variation of reported z_g values in the literature could imply that conventional methods for determining z_g may involve systematic errors in the identification of the critical isotherm (and thus T_g), a problem over which we believe the techniques applied here provide greater control. Alternatively, the field variations of z_g could result from the glass correlation length becoming smaller than the distance between field-induced vortices.^{8,11} Finally, z_g variations could imply an absence of dynamic universality. These possibilities are not investigated in further detail here. However, to study

this issue in more depth, it will be important to apply the methods developed in this work to other high- T_c samples. In each case, universality must be demonstrated for both the scaling exponents and the functions.

Finally, through Eqs. (6) and (7), evidence has been presented which, together with analyses such as Ref. 15, provides strong support for 3D XY multicritical scaling. The present analysis involves scaling of the *nonlinear* resistivity, and therefore complements previous 3D XY studies,^{15,17,24} which apply only to the *linear* resistivity. The focus of the present work has been the vortex glass critical regime, for which Eqs. (3), (6), and (7) augment previous analyses by explaining the field dependence of vortex glass scaling, in addition to the temperature dependence. Through these equations it is found that $z_{xy} \simeq 2.3$, in agreement with Ref. 23; the present result does not suffer from background subtraction difficulties. Eqs. (6) and (7) apply only to the 3D XY scaling region, thus allowing the range of 3D XY scaling to be determined. For the sample studied here, the 3D XY critical regime extends to 8 T along the line $T = T_g(H)$. If the dynamics of the glass transition are ultimately determined to be non-universal, it will be important to ascertain whether the self-consistency of 3D XY multicritical scaling, demonstrated here, still persists.

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FIG. 1. Crossing point method at $H = 1$ T. Different curves correspond to different current densities. The curves all cross at a single point, $T = 65.2$ K, provided z_g is chosen correctly, thus identifying T_g . The z_g values obtained in this way are shown in the inset for each applied field. Additionally, by determining the slopes of the curves at T_g , the exponent ν_g can be obtained, as described in the text. These ν_g values are shown for each applied field in the inset.

FIG. 2. Slope method for $H = 1$ T. The static scaling exponent ν_g is calculated from the power-law dependence of the ohmic $\rho(T)$ data. Values of T_g and $z_g = 3.8$ used in this fit were obtained from the crossing point method of Fig. 1. Resulting values of ν_g are shown in the inset for each applied field.

FIG. 3. Scaled $\rho(J, T)$ data from three different fields. ($\tilde{\rho}$ and \tilde{J} represent, respectively, the left hand side of Eq. (3), and the argument (x) of the scaling function $G_{\pm}[x]$. Dimensionless, arbitrary units are used.) The quantities ρ_0 and J_0 of Eq. (3) have been found for each field so that the three curves collapse onto a single, universal curve. The scaling exponents $z_g = 3.8$ and $\nu_g = 1.8$ are used for

all fields. Although the collapse is excellent for all five fields used here, only three are shown for clarity. The dashed lines represent asymptotic behaviors, and their crossing identifies the normalized point $(1, 1)$; see Eqs. (4) and (5).

FIG. 4. Field dependence of the quantities ρ_0 and J_0 . The dashed lines represent best fits, involving data points only from the expected 3D XY scaling region¹⁵ ($H < 8$ T).